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Extraction of the pion distribution amplitude from polarized muon pair production

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Abstract

We consider the production of muon pairs from the scattering of pions on longitudinally polarized protons. We calculate the cross section and the single spin asymmetry for this process, taking into account pion bound state effects. We work in the kinematic region where the photon has a large longitudinal momentum fraction, which allows us to treat the bound state problem perturbatively. Our predictions are directly proportional to the pion distribution amplitude. A measurement of the polarized Drell-Yan cross section thus allows the determination of the shape of the pion distribution amplitude.

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I. INTRODUCTION

Spin effects are known to provide sensitive tests for the underlying particle theories in many cases. One may recall the spin asymmetries in pion-nucleon scattering which cannot be described in the simplest pole approximation of the Regge theory and require to take into account cuts. In the case of QCD the so-called EMC Spin Crisis was a signal of a breakdown of the naive parton model and of the importance of the incorporation of such a subtle field theory effect as the axial anomaly [1].

A less popular, but also extremely interesting example is the problem of single transverse spin asymmetries in reactions mediated by strong interactions. As these effects require a mass parameter and an imaginary part in the scattering amplitude, they are proportional to $m_q \alpha_s$ in the 'naive' perturbative approach. However, a careful application of QCD factorization at twist-3 level removes both of these small parameters [2]. The collective gluon field of the hadron is shifting the current quark mass m_q to a mass parameter of the order of the hadron mass and, simultaneously, provides the imaginary part. Technically, the latter appears when the pole of the quark (or gluon [3]) propagator is integrated over the light-cone momentum fraction.

A similar imaginary part is present in the QCD higher twist contributions to dilepton production in pion-nucleon scattering, due to the integration over the light-cone momentum fraction of the quark in the pion. This effect leads—in the case of unpolarized nucleon targets— to a significant contribution to the dilepton angular distribution [4], which is sensitive to different ansätze for the pion distribution function.

It is interesting to look for single spin asymmetries for which the imaginary part does not constitute a (large) correction, but the whole contribution. The natural candidate is the spin asymmetry in the production of dileptons from the scattering of pions on a *longitudinally polarized* nucleon. This spin asymmetry has been thoroughly studied in perturbative QCD [5,6]. The analogous mesonic processes have been discussed in [7]. In the present work we study this single spin asymmetry, taking into account pion bound state effects.

II. CALCULATION OF THE HADRONIC TENSOR

We will consider the inclusive production of muon pairs from the scattering of a pion on a polarized proton (the Drell-Yan process),

$$\pi^-(P_\pi) + p(P_p, s_\ell) \to \gamma^*(Q,\epsilon) + X \to \mu^+(q_+) + \mu^-(q_-) + X, \tag{2.1}$$

where s_ℓ denotes the degree of longitudinal polarization of the proton, and all momenta refer to the overall c.m. system. It is convenient to write the differential cross section for (2.1) in terms of a hadronic and a leptonic tensor,

$$d\sigma = \frac{1}{2s}\left(\frac{1}{Q^2}\right)^2 \frac{d^3q_+}{(2\pi)^3 2q_+^0}\frac{d^3q_-}{(2\pi)^3 2q_-^0} W^{\mu\nu}(s_\ell) L_{\mu\nu}, \tag{2.2}$$

where s is the hadronic c.m. energy. The leptonic tensor has the well-known form

$$L^{\mu\nu} = 4e^2\left(q_+^\mu q_-^\nu + q_-^\mu q_+^\nu - g^{\mu\nu}Q^2/2\right). \tag{2.3}$$

Note that $L^{\mu\nu}$ is symmetric, since we only consider the exchange of a virtual photon (not a Z boson) and since we sum over the polarizations of $\mu^\pm$.

The hadronic tensor $W^{\mu\nu}$ may be written as

$$\begin{aligned} W^{\mu\nu}(s_\ell) = & \sum_X (2\pi)^4 \delta^4(Q + P_X - P_\pi - P_p) \\ & \langle \gamma(Q,\epsilon^\mu), X(P_X)|\mathcal{T}|\pi^-(P_\pi), p(P_p, s_\ell)\rangle \\ & \langle \gamma(Q,\epsilon^\nu), X(P_X)|\mathcal{T}|\pi^-(P_\pi), p(P_p, s_\ell)\rangle^*. \end{aligned} \tag{2.4}$$

Only the symmetric part of $W^{\mu\nu}$ is relevant after contraction with $L_{\mu\nu}$ of eq. (2.3). Therefore $W^{\mu\nu}$ in the following is implicitly assumed to be symmetrized.

The angular distribution of μ^+ in (2.1) may be parametrized as

$$\begin{aligned} \frac{d\sigma}{dQ^2 dQ_T^2 dx_L d\cos\theta d\phi} \propto & \; 1 + \lambda\cos^2\theta + \mu\sin 2\theta\cos\phi + \frac{\nu}{2}\sin^2\theta\cos 2\phi \\ & + \bar{\mu}\sin 2\theta\sin\phi + \frac{\bar{\nu}}{2}\sin^2\theta\sin 2\phi. \end{aligned} \tag{2.5}$$

Here, θ and ϕ are angles defined in the muon pair rest frame, and λ, μ, ν, $\bar{\mu}$, $\bar{\nu}$ are angle independent coefficients. They depend on Q^2 (the virtuality of the photon), Q_T^2 (the

squared transverse momentum of the photon in the hadronic c.m.s.), and on $x_L = 2Q_L/\sqrt{s}$ (the longitudinal momentum fraction of the photon). We take $x_L > 0$ for a photon moving forward with respect to the pion in the hadronic c.m. system. The normalization of the cross section can be determined from

$$\frac{Q^2 d\sigma}{dQ^2 dQ_T^2 dx_L} = \frac{e^2}{96(2\pi)^4} \frac{(3+\lambda)N}{\sqrt{s}\sqrt{Q^2 + Q_T^2 + x_L^2 s/4}}. \tag{2.6}$$

The normalization N and the angular coefficients are related to the helicity amplitudes of the hadronic tensor:

$$\begin{aligned}
N &= W_T + W_L = \epsilon_\mu(+1) W^{\mu\nu} \epsilon_\nu^*(+1) + \epsilon_\mu(0) W^{\mu\nu} \epsilon_\nu(0) \\
\lambda &= N^{-1}(W_T - W_L) = N^{-1}\{\epsilon_\mu(+1) W^{\mu\nu} \epsilon_\nu^*(+1) - \epsilon_\mu(0) W^{\mu\nu} \epsilon_\nu(0)\} \\
\mu &= N^{-1} W_{LT} = (\sqrt{2}N)^{-1}\{\epsilon_\mu(0) W^{\mu\nu} \epsilon_\nu(+1) + \epsilon_\mu(0) W^{\mu\nu} \epsilon_\nu^*(+1)\} \\
\nu &= N^{-1} W_{TT} = N^{-1}\{\epsilon_\mu(+1) W^{\mu\nu} \epsilon_\nu^*(-1) + \epsilon_\mu^*(+1) W^{\mu\nu} \epsilon_\nu(-1)\} \\
\bar{\mu} &= N^{-1} \overline{W}_{LT} = i(\sqrt{2}N)^{-1}\{\epsilon_\mu(0) W^{\mu\nu} \epsilon_\nu(+1) - \epsilon_\mu(0) W^{\mu\nu} \epsilon_\nu^*(+1)\} \\
\bar{\nu} &= N^{-1} \overline{W}_{TT} = i N^{-1}\{\epsilon_\mu(+1) W^{\mu\nu} \epsilon_\nu^*(-1) - \epsilon_\mu^*(+1) W^{\mu\nu} \epsilon_\nu(-1)\},
\end{aligned} \tag{2.7}$$

where the polarization vectors $\epsilon^\mu(\pm 1) = (0, \mp \mathbf{e}_x + i\mathbf{e}_y)$ and $\epsilon^\mu(0) = (0, \mathbf{e}_z)$ are determined by specifying the coordinate axes $\mathbf{e}_i$ $(i = x, y, z)$ in the muon rest frame.

The coefficients $\bar{\mu}$ and $\bar{\nu}$ are nonzero only in the polarized Drell-Yan process and are induced by absorptive parts in the scattering amplitude $\mathcal{T}$. In the parton model, the leading contributions to $\bar{\mu}$ and $\bar{\nu}$ come from the interference of Born diagrams with the absorptive parts of one-loop diagrams [6]. The relation of the single spin asymmetry

$$\mathcal{A} \equiv \frac{d\sigma(s_\ell = +1) - d\sigma(s_\ell = -1)}{d\sigma(s_\ell = +1) + d\sigma(s_\ell = -1)} \tag{2.8}$$

to the angular coefficients is simply given by

$$\mathcal{A} = \frac{\bar{\mu}(s_\ell = +1)\sin 2\theta \sin\phi + \frac{1}{2}\bar{\nu}(s_\ell = +1)\sin^2\theta \sin 2\phi}{1 + \lambda\cos^2\theta + \mu \sin 2\theta \cos\phi + \frac{1}{2}\nu \sin^2\theta \cos 2\phi}. \tag{2.9}$$

For the unpolarized Drell-Yan process the parton model predictions [8] fail to describe the data on λ, μ and ν [9]. In contrast, taking into account pion bound state effects

yields reasonable fits to the data [4]. An interesting feature of this bound state model is a nonvanishing absorptive part in *leading* order which is proportional to the pion distribution amplitude φ. This property directly leads to nonvanishing $\bar{\mu}$ and $\bar{\nu}$ which are strongly sensitive to φ. In the following we will shortly describe the salient features of the bound state model for the pion. Further details may be found, e.g., in [10], [4].

The bound state problem may be treated perturbatively if the momentum fraction of the quark from the pion is large, which is the case for large x_L. The dominant contribution to the Drell-Yan process comes from the annihilation of a quark with an antiquark [11]. In the region of large x_L, the diagrams of fig. 1a,b. then give the leading contributions. We will explain the formalism for the annihilation of a u quark from the proton with the $\bar{u}$ quark from the pion; the other possible annihilation process of a $\bar{d}$ quark from the proton with the d quark from the pion is treated in complete analogy and its contribution will later be added incoherently.

In diagram 1a we see that the $\bar{u}$ quark propagator is far off-shell, $p_{\bar{u}}^2 = -Q_T^2/(1 - x_{\bar{u}})$, if the light cone momentum fraction of the quark from the pion $x_{\bar{u}} \approx x_L$ is close to 1 [10]. Thus, the pion can be resolved by a single hard gluon exchange [12]. The second diagram, fig. 1b, is required by gauge invariance.

The hadronic tensor $W^{\mu\nu}$ is obtained by a convolution of the partonic tensor $w^{\mu\nu}$ with the corresponding parton distribution functions for the polarized proton[1]. The partonic tensor $w^{\mu\nu}$ is computed from the product $M^\mu M^{*\nu}$. M^μ is the amplitude for the reaction

$$u^\uparrow + \pi^- \to \gamma^*(Q, \epsilon^\mu) + X, \tag{2.10}$$

where the arrow indicates the polarization of the parton. This amplitude M^μ is obtained by convoluting the partonic amplitude $T^\mu(u^\uparrow + \bar{u}d \to \gamma^* + d)$ with the pion distribution amplitude $\varphi(z, \tilde{Q}^2)$ [12],

---

[1]Here we take into account the correct partonic flux factor which amounts to multiplying with $s/\hat{s}$, $\hat{s} = (P_\pi + p_u)^2$.

$$M^\mu = \frac{f_\pi}{4N_c} \int_0^1 dz\ \varphi(z, \tilde{Q}^2)\ T^\mu, \qquad \int_0^1 dz\ \varphi(z, \tilde{Q}^2) = 1, \tag{2.11}$$

where $\tilde{Q}^2 \sim Q_T^2/(1-x_{\bar{u}})$ is the cutoff for the integration over soft momenta in the definition of φ, $f_\pi \approx 133$ MeV is the pion decay constant, and $N_c = 3$ is the number of colors. A detailed discussion of the scale setting problem for exclusive processes is given in [13]. Using the commensurate scale relations in the calculation of (2.11) would allow to eliminate ambiguities caused by the factorization and renormalization. However, this does not change the qualitative features of our predictions and thus it will not be discussed in this paper.

The imaginary part in this model arises because the internal quark line of fig. 1b can go on-shell. After integrating the hard scattering amplitude T^μ according to equation (2.11), we are left with a regular amplitude M^μ which has an imaginary part proportional to the pion distribution amplitude at the point where the quark propagator gets singular.

We are now ready to present our analytic results. We relegate the full result for the hadronic tensor to the appendix, and confine ourselves here to a discussion of the angular coefficients $\bar{\mu}$ and $\bar{\nu}$. To present our results for the angular distribution we choose the Gottfried-Jackson frame where the $\mathbf{e}_z$ axis is taken to be the pion direction in the muon rest frame and the $\mathbf{e}_x$ axis lies in the $\pi^- P$ plane such that the proton momentum has a negative x component. From eqs. (2.7), (A1), and (A2) we get the angular coefficients

$$\begin{aligned} \bar{\mu} = &\frac{-2 s_\ell \rho \tilde{x} F \pi \varphi(\tilde{x})}{(1-\tilde{x})^2 \left[(F + \mathrm{Re}\ I(\tilde{x}))^2 + \pi^2 \varphi(\tilde{x})^2\right] + \rho^2 \tilde{x}^2 F^2 (4+\rho^2)} \\ &\times \frac{\frac{4}{9}\Delta q_u^v(x_p) + \frac{4}{9}\Delta q_u^s(x_p) + \frac{1}{9}\Delta q_d^s(x_p)}{\frac{4}{9} q_u^v(x_p) + \frac{4}{9} q_u^s(x_p) + \frac{1}{9} q_d^s(x_p)}, \\ \bar{\nu} = &\ 2\rho\bar{\mu}. \end{aligned} \tag{2.12}$$

Here, $\rho = Q_T/Q$ and $\tilde{x}$ is a function of $x_L, Q^2/s$ and ρ

$$\tilde{x} \equiv \frac{x_{\bar{u}}}{1+\rho^2} = \frac{1}{2} \frac{x_L + \sqrt{x_L^2 + 4Q^2 s^{-1}(1+\rho^2)}}{1+\rho^2}. \tag{2.13}$$

Furthermore,

$$\begin{aligned} F &= \int_0^1 dz \frac{\varphi(z, \tilde{Q}^2)}{z}, \\ I(\tilde{x}) &= \int_0^1 dz \frac{\varphi(z, \tilde{Q}^2)}{z(z + \tilde{x} - 1 + i\epsilon)} \end{aligned} \tag{2.14}$$

denote integrals over the pion distribution amplitude. Finally, $q_i^j(x_p)$ and $\Delta q_i^j(x_p)$ are the unpolarized and polarized valence and sea quark distribution functions evaluated at the point $x_p \approx Q^2/(s\tilde{x})$. The results for λ, μ, and ν within the pion bound state model are given elsewhere [4].

III. NUMERICAL RESULTS

The numerical results for the angular coefficient functions $\bar{\mu}$ and $\bar{\nu}$, eq. (2.12), depend on $\tilde{x}$, ρ, the pion distribution amplitude and on the ratio of the polarized and unpolarized quark distribution functions of the proton. Whereas the unpolarized quark distributions are known quite well, the polarized ones still contain large uncertainties. For our results we choose the parametrizations given in [14] at $Q^2 = 4$ GeV^2.

In fig. 2 we present our results for several distribution amplitudes φ that are quite different in shape. Fig. 2a shows our choices for φ: The solid line represents the two-humped function [15] which gave a good fit to the data on λ, μ and ν [4]. The effective evolution parameter is set to $\tilde{Q}^2 \sim 4$ GeV^2. Since the predictions are also very sensitive to the behavior of the distribution amplitude in the endpoint region, we choose two different parametrizations [16] for the convex distribution amplitude $\varphi(z) = z^a(1-z)^a/\mathrm{B}(a+1, a+1)$. The dashed line in fig. 2a represents the asymptotical amplitude, i.e., $a = 1$, and the dotted line shows a narrow distribution amplitude with $a = 10$.

In fig. 2b the moment $\int \sin 2\theta \sin\phi d\sigma(s_\ell = +1) \propto N\bar{\mu}$ is plotted in arbitrary units versus x_L for the values $s = 20^2$ GeV^2, $Q = 3$ GeV and $Q_T = 0.9$ GeV. It can be seen that different shapes of the distribution amplitude could be distinguished by a measurement of this quantity, which is proportional to the pion distribution amplitude evaluated at $\tilde{x} \approx x_L$ (cf. the numerator of (2.12)). It is also demonstrated that narrow distribution amplitudes give a moment which vanishes in the large x_L region.

In figs. 2c and 2d the angular coefficient $\bar{\mu}(s_\ell = +1)$ is plotted versus x_L for two different choices of Q_T and the same values for s and Q as in fig. 2b. In fig. 2c Q_T was set again to

0.9 GeV, which gives the moderate value $\rho = 0.3$. It is demonstrated that the two-humped form for φ induces a minimum at $x_L \sim 0.6$ which, however, vanishes for the smaller value $\rho = 0.06$ used in fig. 2d. For very narrow distribution amplitudes $\bar{\mu}$ is strongly suppressed. Since $\bar{\nu} = 2\rho\bar{\mu}$ (cf. eq. (2.12)) in the bound state model, we do not show separate plots for this quantity.

Note that the effects of a variation of the pion distribution amplitude in the spin-dependent and spin-averaged cross sections – i.e. in the numerator and denominator of (2.12) – partially compensate each other. As a consequence, the result for the two-humped distribution amplitude appears in between the results for the two convex distributions. The spin-dependent part of the cross section is a more sensitive "partonometer", than the dimensionless coefficients $\bar{\mu}, \bar{\nu}$. Although the former quantity suffers from larger uncertainties due to higher order corrections, , one may expect only a small uncertainty for the determination of the shape of pion distribution amplitude, if the x_L dependence of the corrections is weak.

Although the naive parton model is at variance with the unpolarized angular distribution [9], we would like to present for completeness and comparison also the predictions of this model for $\bar{\mu}$ and $\bar{\nu}$. We took the analytic results from [6] and again used the parton distribution functions of [14]. In the range $0.5 < x_L < 1$ the parton model yields coefficients $\bar{\mu}(s_\ell = +1)$ and $\bar{\nu}(s_\ell = +1)$ which are positive and to a very good approximation independent of x_L. For the same values used in fig2b and 2c, i.e. $s = 20^2$ GeV$^2, Q = 3$ GeV and $\rho = 0.3$ we get a value of $\bar{\mu}(s_\ell = +1) \sim 0.12\alpha_s \sim 0.036$. This means that bound state effects are roughly of the same magnitude but opposite in sign compared to the parton model at moderate values of ρ. For the smaller value $\rho = 0.06$ we find $\bar{\mu}(s_\ell = +1) \sim 0.034\alpha_s \sim 0.01$, again independent of x_L and positive. At such a small value of ρ the bound state model predicts that $\bar{\mu}$ and $\bar{\nu}$ are concentrated at large x_L for not too narrow distribution amplitudes. In the kinematical range considered, the relation between $\bar{\nu}$ and $\bar{\mu}$ calculated within the parton model takes a similar form as in the bound state model, namely, $\bar{\nu} \sim 3\rho\bar{\mu}$.

In conclusion we have shown that experiments on dimuon production from the scattering of pions on polarized targets will give detailed information both on the shape and on the

endpoint behavior of the pion distribution amplitude.

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APPENDIX: RESULT FOR THE HADRONIC TENSOR

Using the kinematic decomposition of the (symmetric part of the) hadronic tensor

$$
\begin{aligned}
W^{\mu\nu} = & -\tilde{g}^{\mu\nu} W_1 + \frac{\tilde{P}_\pi^\mu \tilde{P}_\pi^\nu}{Q^2} W_2 + \frac{\tilde{P}_p^\mu \tilde{P}_p^\nu}{Q^2} W_3 + \frac{(\tilde{P}_\pi^\mu \tilde{P}_p^\nu + \tilde{P}_p^\mu \tilde{P}_\pi^\nu)}{Q^2} W_4 \\
& + \frac{\left(\epsilon^\mu{}_{\alpha\beta\gamma} \tilde{P}_\pi^\nu + \epsilon^\nu{}_{\alpha\beta\gamma} \tilde{P}_\pi^\mu\right) P_\pi^\alpha P_p^\beta Q^\gamma}{Q^4} s_\ell W_5 \\
& + \frac{\left(\epsilon^\mu{}_{\alpha\beta\gamma} \tilde{P}_p^\nu + \epsilon^\nu{}_{\alpha\beta\gamma} \tilde{P}_p^\mu\right) P_\pi^\alpha P_p^\beta Q^\gamma}{Q^4} s_\ell W_6, \qquad \text{(A1)}
\end{aligned}
$$

where $\tilde{g}^{\mu\nu} = g^{\mu\nu} - Q^\mu Q^\nu / Q^2$ and $\tilde{P}_i^\mu = \tilde{g}^\mu{}_\nu P_i^\nu$, we obtain from the bound state model the following predictions for the structure functions:

$$
\begin{aligned}
W_1 &= \frac{n}{4} \left\{ \left(\left[1 - \tilde{x}(1+\rho^2)\right] F + (1-\tilde{x}) \mathrm{Re}\, I(\tilde{x}) \right)^2 + \pi^2 \varphi(\tilde{x})^2 \right\} f(x_p) \\
W_2 &= -n \tilde{x}^2 f(x_p) \\
& \qquad \left\{ (F + (1-\tilde{x}) \mathrm{Re}\, I(\tilde{x})) \left[(1 - \tilde{x}(1+\rho^2)^2) F + (1-\tilde{x}) \mathrm{Re}\, I(\tilde{x})) \right] + \pi^2 \varphi(\tilde{x})^2 \right\} \\
W_3 &= \frac{n x_p^2 \tilde{x}}{1-\tilde{x}} \left[1 - \tilde{x}(1+\rho^2) \right]^2 F(F + \mathrm{Re}\, I(\tilde{x})) f(x_p) \\
W_4 &= -\frac{n x_p \tilde{x}^2}{1-\tilde{x}} \rho^2 \left[1 - \tilde{x}(1+\rho^2) \right] F(F + (1-\tilde{x}) \mathrm{Re}\, I(\tilde{x})) f(x_p)
\end{aligned}
$$

$$
\begin{aligned}
W_5 &= -\frac{n x_p \tilde{x}^3}{1-\tilde{x}}(1+\rho^2)\left[1-\tilde{x}(1+\rho^2)\right]\pi\varphi(\tilde{x})F\Delta f\,(x_p) \\
W_6 &= \frac{n x_p^2 \tilde{x}^2}{(1-\tilde{x})^2}\left[1-\tilde{x}(1+\rho^2)\right]^2\pi\varphi(\tilde{x})F\Delta f\,(x_p)\,. \qquad \text{(A2)}
\end{aligned}
$$

Here,

$$
n = \frac{32\pi^3 e^2 \alpha_s^2 C_F^2 f_\pi^2}{N_c^2 Q_T^4} \quad \text{with } C_F = \frac{N_c^2-1}{2N_c} \qquad \text{(A3)}
$$

is a normalization factor, and

$$
\begin{aligned}
f(x_p) &= \frac{4}{9}q_u^v(x_p) + \frac{4}{9}q_u^s(x_p) + \frac{1}{9}q_d^s(x_p) \\
\Delta f(x_p) &= \frac{4}{9}\Delta q_u^v(x_p) + \frac{4}{9}\Delta q_u^s(x_p) + \frac{1}{9}\Delta q_d^s(x_p), \qquad \text{(A4)}
\end{aligned}
$$

are given by the unpolarized and polarized valence and sea quark distribution functions. The variable $\tilde{x}$ and x_p as well as the integrals F and Re I containing the pion distribution amplitude are defined below eq. (2.12).

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FIGURES

FIG. 1. Leading contributions to the amplitude for the Drell-Yan process at large x_L.

FIG. 2. In (a) we show different pion distribution amplitudes $\varphi(z)$. Solid line: two-humped function with an effective evolution parameter of $\tilde{Q}^2 \sim 4$ GeV^2; dashed line: $\varphi(z) = z^a(1-z)^a/\mathrm{B}(a+1,a+1)$ with $a=1$; dotted line: same with $a=10$. In (b) the moment $\int \sin 2\theta \sin\phi d\sigma(s_\ell = 1)$ is plotted versus x_L for the values $s = 20^2$ GeV^2, $Q = 3$ GeV and $Q_T = 0.9$ GeV. In (c) and (d) the angular coefficient $\bar{\mu}(s_\ell = 1)$ is plotted versus x_L for the two different values $Q_T/Q = 0.3$ and $Q_T/Q = 0.06$, respectively, and for the same values for s and Q as in (b). In (b-d) we use the same choices for $\varphi(z)$ as in (a).

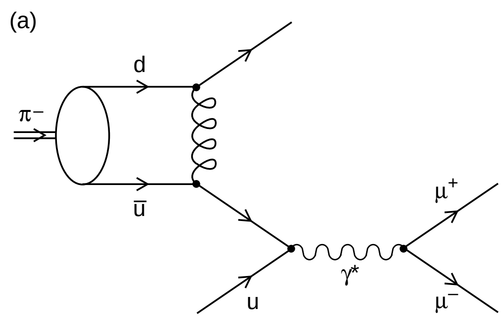

Figure 1

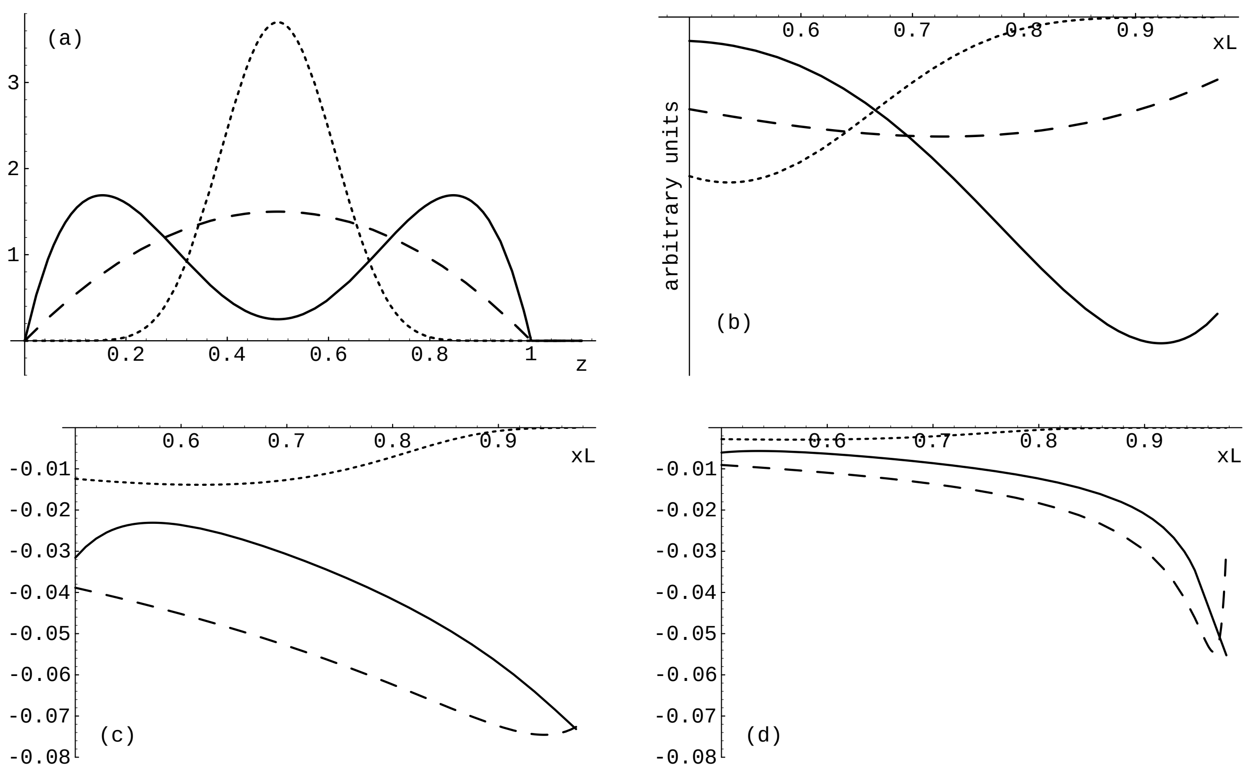

Figure 2